

The $\mu\nu$ SSM with an extra $U(1)$

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ABSTRACT: The $\mu\nu$ SSM solves the μ problem of the MSSM and generates correct neutrino masses by simply using right-handed neutrinos. This mechanism implies that only dimensionless trilinear terms, breaking R -parity, are present in the superpotential. We present an extension of the $\mu\nu$ SSM with an extra $U(1)$ gauge symmetry. We use the extra $U(1)$ charges of the matter fields to forbid the presence in the superpotential of renormalizable and non-renormalizable baryon number violating operators, the trilinear operator producing a domain wall problem, and bilinear operators such as the μ term and the Majorana masses. We apply the anomaly cancellation conditions associated to the extra $U(1)$, to constrain the values of the $U(1)$ charges. We find that six assignments of the $U(1)$ charges to the matter fields are viable, once extra matter is introduced. In particular, three generations of vector-like color triplets and $SU(2)_L$ doublets, as well as six Standard Model singlets are necessary. Electroweak symmetry breaking is viable in the model, with wide regions of the parameter space fulfilling the experimental constraints on the existence of a new gauge boson Z' . Neutrinos and the extra gaugino mix with the MSSM neutralinos, producing a generalized see-saw matrix that can reproduce the experimental results on neutrino masses. Finally, we have estimated the tree-level upper bound on the lightest Higgs mass, finding that it can be as large as about 120 GeV.

KEYWORDS: Supersymmetric Effective Theories, Beyond Standard Model, Supersymmetry Phenomenology.

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1. Introduction

The μ from ν Supersymmetric Standard Model ($\mu\nu$ SSM) [1, 2, 3] is defined by the following superpotential:

$$W = \epsilon_{ab} \left(Y_u^{ij} \hat{H}_2^b \hat{Q}_i^a \hat{u}_j^c + Y_d^{ij} \hat{H}_1^a \hat{Q}_i^b \hat{d}_j^c + Y_e^{ij} \hat{H}_1^a \hat{L}_i^b \hat{e}_j^c + Y_\nu^{ij} \hat{H}_2^b \hat{L}_i^a \hat{\nu}_j^c \right) - \epsilon_{ab} \lambda^i \hat{\nu}_i^c \hat{H}_1^a \hat{H}_2^b + \frac{1}{3} \kappa^{ijk} \hat{\nu}_i^c \hat{\nu}_j^c \hat{\nu}_k^c, \quad (1.1)$$

where we take $\hat{H}_1^T = (\hat{H}_1^0, \hat{H}_1^-)$, $\hat{H}_2^T = (\hat{H}_2^+, \hat{H}_2^0)$, $\hat{Q}_i^T = (\hat{u}_i, \hat{d}_i)$, $\hat{L}_i^T = (\hat{\nu}_i, \hat{e}_i)$, $i, j, k = 1, 2, 3$ and $a, b = 1, 2$ are generation and $SU(2)$ indices, respectively, and $\epsilon_{12} = 1$. This superpotential contains in addition to the usual Yukawas for quarks and charged leptons, Yukawas for neutrinos $Y_\nu \hat{H}_2 \hat{L} \hat{\nu}^c$, terms of the type $\lambda \hat{\nu}^c \hat{H}_1 \hat{H}_2$ producing an effective μ term through right-handed sneutrino vacuum expectation values (VEVs) of the order of the electroweak (EW) scale, $\mu \equiv \lambda \langle \tilde{\nu}^c \rangle$, and terms of the type $\kappa \hat{\nu}^c \hat{\nu}^c \hat{\nu}^c$ avoiding the existence of a Goldstone boson and producing an EW-scale see-saw through the generation of effective Majorana masses $\kappa \langle \tilde{\nu}^c \rangle$.

Thus the $\mu\nu$ SSM solves the μ -problem [4] of the Minimal Supersymmetric Standard Model (MSSM) [5] and generates light neutrino masses by simply using right-handed neutrino superfields. Note that the above terms in the superpotential produce the explicit breaking of R -parity in this model. The size of the breaking can be easily understood realizing that in the limit where Y_ν are vanishing, the $\hat{\nu}^c$ are ordinary singlet superfields like the \hat{S} of the Next-to-Minimal Supersymmetric Standard Model (NMSSM) [6], without any connection with neutrinos, and R -parity is therefore conserved. Once Y_ν are switched on, the $\hat{\nu}^c$ become right-handed neutrinos, and, as a consequence, R -parity is broken. Thus the breaking is small because the EW-scale see-saw implies small values for $Y_\nu \sim 10^{-6}$.

Since the $\mu\nu$ SSM is a very well motivated and attractive model, several phenomenological studies have been carried out. In [7, 8], the parameter space, the spectrum and the vacua of the $\mu\nu$ SSM were analyzed in detail. The neutrino sector was studied in [7, 8, 9, 10],

obtaining that current neutrino data (the measured mass differences and mixing angles) can be easily reproduced. Analyses of possible signals at colliders were also carried out. Since the Lightest Supersymmetric Particle (LSP) is no longer stable due to the breaking of R -parity, not all supersymmetric chains must yield missing energy events. In [9, 11, 12] the decays of the lightest neutralino were discussed, as well as the correlations of the decay branching ratios with the neutrino mixing angles. Let us remark that the breaking of R -parity generates a peculiar structure for the mass matrices, and this has to be taken into account in the computations mentioned above. In particular, the presence of right and left-handed sneutrino VEVs leads to mixing of neutralinos with left- and right-handed neutrinos, and as a consequence a generalized matrix of the see-saw type. Besides, there is also the mixing of the neutral Higgses with the sneutrinos producing 8×8 neutral scalar mass matrices, and this extended Higgs sector could be very helpful for testing the $\mu\nu$ SSM at colliders [11, 13]. In [14], special emphasis was put in the decays of the Higgses and viable benchmark points for LHC searches were provided.

Concerning cosmological issues, dark matter and baryon asymmetry have been analyzed in the model. The gravitino, present in the local supersymmetric version of the model, could be a good dark matter candidate as discussed in [15, 16], where its possible detection through the observation of a monochromatic gamma-ray line in the *Fermi* satellite was also studied. In [17], the generation of the baryon asymmetry of the Universe was analyzed in detail in the context of the $\mu\nu$ SSM, with the interesting result that EW baryogenesis can be realized.

Once R -parity is not a symmetry of the model, lepton and baryon number violating terms in the superpotential like

$$\epsilon_{ab} \left(\lambda'''_{ijk} \hat{L}_i^a \hat{L}_j^b \hat{e}_k^c + \lambda'_{ijk} \hat{L}_i^a \hat{Q}_j^b \hat{d}_k^c + \mu_i \hat{L}_i^a \hat{H}_2^b \right) , \quad \lambda''_{ijk} \hat{d}_i^c \hat{d}_j^c \hat{u}_k^c , \quad (1.2)$$

are in principle allowed by gauge invariance. As it is well known, to avoid too fast proton decay mediated by the exchange of squarks of masses of the order of the EW scale, the presence together of terms of the type $\hat{L}\hat{Q}\hat{d}^c$ and $\hat{d}^c\hat{d}^c\hat{u}^c$ must be forbidden, unless we impose very stringent bounds such as e.g. $\lambda_{112}^{*'}\lambda''_{112} \lesssim 2 \times 10^{-27}$. Clearly, these values for the couplings are not very natural, and for constructing viable supersymmetric models one usually forbids at least one of the operators LQd^c or $d^c d^c u^c$. The other type of operators above are not so stringently suppressed, and therefore still a lot of freedom remains [18].

There are several ways to avoid this problem. One is to assume that there are other discrete symmetries, like e.g. baryon triality which only forbids the baryon violating operators [19]. Another one comes from string constructions, where the matter superfields can be located in different sectors or have different extra $U(1)$ charges, in such a way that some operators violating R -parity can be forbidden [20], but others can be allowed.

Another problem is related to the absence of a μ term as well as Majorana masses for neutrinos in the superpotential (1.1), since both type of bilinear terms are in principle allowed by gauge invariance. As for the proton decay problem above, we have several solutions. The fact that only dimensionless trilinear terms are present in the superpotential of the $\mu\nu$ SSM, can be explained invoking a Z_3 symmetry, as it is usually done in the NMSSM. The second solution comes again from string constructions, where the low-energy

limit is determined by the massless string modes. Since the massive modes are of the order of the string scale, only trilinear couplings are present in the low-energy superpotential.

Finally, since the superpotential of the $\mu\nu$ SSM contains only trilinear couplings, it has a Z_3 symmetry, just like the NMSSM, as mentioned above. Therefore, one expects to have also a cosmological domain wall problem [21, 22] in this model. Nevertheless, the usual solution [23] can also work in this case: non-renormalizable operators [21] in the superpotential can explicitly break the dangerous Z_3 symmetry, lifting the degeneracy of the three original vacua, and this can be done without introducing hierarchy problems. In addition, these operators can be chosen small enough as not to alter the low-energy phenomenology.

The aim of this work is to solve the above three problems adopting a different strategy. In particular, we will add an extra $U(1)$ gauge symmetry to the gauge group of the Standard Model. In this way, and since all the fields of the $\mu\nu$ SSM can be charged under the extra $U(1)$, all the dangerous operators could be forbidden without relying in string theory arguments, discrete symmetries or non-renormalizable operators. Previous works using an extra $U(1)$ to solve these problems in other models, can be found in [24, 25].

The outline of the paper is as follows. In Section 2, first we will use the extra $U(1)$ charges of the matter fields to allow the presence of the phenomenologically interesting operators, forbidding the dangerous ones. Then, we will impose the anomaly cancellation conditions associated to the extra $U(1)$ to constrain the values of the $U(1)$ charges. We will see that several assignments of the $U(1)$ charges to the matter fields are viable, but in all cases the introduction of extra matter is required. Once we have found consistent assignments (models), in Section 3 we will study their phenomenology concerning the EW symmetry breaking. We will also check that the experimental constraints on the existence of an extra gauge boson are fulfilled, as well as that correct neutrino masses can be obtained. The tree-level upper bound on the lightest Higg boson mass will also be discussed. Finally, the conclusions are left for Section 4.

2. The search of models

As mentioned in the Introduction, we will work with the gauge group of the Standard Model adding an extra $U(1)$,

$$SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{extra} . \quad (2.1)$$

The matter content of the $\mu\nu$ SSM with three families of quarks and leptons and one family of Higgses has then the following representations under this gauge group:

$$\begin{aligned} & Q(3, 2, \frac{1}{6}, Q_Q) , \quad u^c(\bar{3}, 1, -\frac{2}{3}, Q_u) , \quad d^c(\bar{3}, 1, \frac{1}{3}, Q_d) , \\ & L(1, 2, -\frac{1}{2}, Q_L) , \quad e^c(1, 1, 1, Q_e) , \quad \nu^c(1, 1, 0, Q_{\nu^c}) , \\ & H_1(1, 2, -\frac{1}{2}, Q_{H_1}) , \quad H_2(1, 2, \frac{1}{2}, Q_{H_2}) , \end{aligned} \quad (2.2)$$

where for simplicity we have taken the extra charges as family independent.

Now we ask the Yukawa terms, $\hat{Q}\hat{H}_1\hat{d}^c, \hat{Q}\hat{H}_2\hat{u}^c, \hat{L}\hat{H}_1\hat{e}^c, \hat{L}\hat{H}_2\hat{\nu}^c$ (that give tree-level masses to all fermions), and the effective μ term, $\hat{\nu}^c\hat{H}_1\hat{H}_2$, to be allowed in the superpotential. Since they have to be invariant under $U(1)_{extra}$, we can obtain five equations for the extra $U(1)$ charges. Using these equations we can express five charges in terms of the other three:

$$Q_u = Q_{H_1} + Q_d - Q_{H_2} , \quad (2.3)$$

$$Q_Q = -Q_{H_1} - Q_d , \quad (2.4)$$

$$Q_e = -2Q_{H_1} , \quad (2.5)$$

$$Q_L = Q_{H_1} , \quad (2.6)$$

$$Q_{\nu^c} = -Q_{H_1} - Q_{H_2} . \quad (2.7)$$

It is worth noticing here that equations (2.4), (2.5) and (2.6), imply that the lepton number violating terms, $\hat{L}\hat{L}\hat{e}^c$ and $\hat{L}\hat{Q}\hat{d}^c$, are automatically allowed. Thus to avoid fast proton decay, the baryon number violating term, $\hat{d}^c\hat{d}^c\hat{u}^c$, should be forbidden. Using (2.3) one obtains the following condition:

$$Q_{H_1} \neq Q_{H_2} - 3Q_d . \quad (2.8)$$

Besides, to forbid the bilinear μ -term, $\mu\hat{H}_1\hat{H}_2$, one has to impose,

$$Q_{H_1} \neq -Q_{H_2} . \quad (2.9)$$

Given (2.6), this implies that the lepton number violating operator $\hat{L}\hat{H}_2$ is automatically forbidden. In addition, from (2.7) one obtains that $Q_{\nu^c} \neq 0$, and, as a consequence, the term that generates the cosmological domain wall problem, $\hat{\nu}^c\hat{\nu}^c\hat{\nu}^c$, is also automatically forbidden. It is worth noticing here that a Goldstone boson does not appear from the absence of this term in the superpotential, since the $U(1)$ symmetry is gauged. As a consequence, the Goldstone boson is eaten by the Z' in the process of EW symmetry breaking. We will see in the next section that the analysis of the generalized see-saw matrix mixing neutrinos with neutralinos, in the case of three generations, is different from the usual one in the $\mu\nu$ SSM because of the absence of this effective Majorana mass term.

Let us now impose the six anomaly cancellation conditions associated to the extra $U(1)$ gauge symmetry.

The cancellation of the anomaly $[SU(2)_L]^2 - U(1)_{extra}$ implies that the condition $\sum Q_{extra} = 0$ must be fulfilled, where the sum extends over all left-handed fermions and antifermions, and Q_{extra} generically denotes the $U(1)_{extra}$ charges of the particles (in this case only doublet fermions). Using in addition (2.4) and (2.6), one obtains for Q_d the solution $Q_d = \frac{n_H-6}{9}Q_{H_1} + \frac{n_H}{9}Q_{H_2}$, where we have assumed in principle that the number of Higgs doublets n_H is free. Note however that if $n_H = 3$ one would obtain $Q_{H_1} = Q_{H_2} - 3Q_d$, which does not fulfill (2.8) and therefore the baryon number violating operator $d^c d^c u^c$ would be allowed. Thus we will not consider this possibility and impose $n_H \neq 3$.

The $[SU(3)_C]^2 - U(1)_{extra}$ anomaly cancellation condition, $\sum Q_{extra} = 0$ (only for color triplet fermions), gives rise to: $3(2Q_Q + Q_u + Q_d) = 0$. But once (2.3) and (2.4) are

taken into account, one obtains $Q_{H_1} = -Q_{H_2}$, which does not fulfill (2.9), thus the bilinear operators would be allowed in the superpotential, spoiling the solution of the $\mu\nu$ SSM to the μ problem. Then, we conclude that *we have to introduce exotic matter with color charge* in the spectrum to cancel this anomaly. On the other hand, in order not to alter the anomaly cancellation conditions associated to the Standard Model gauge group, we assume that we have n_q generations of exotics which are vector-like pairs of chiral superfields with opposite $U(1)_Y$ hypercharges: $\hat{q}(3, 1, Y_q, Q_q)$, $\hat{q}^c(\bar{3}, 1, -Y_q, Q_{q^c})$. In addition, to avoid conflicts with experiments, the exotic quarks must be sufficiently heavy to not have been detected. Thus to give them masses, we add a trilinear term in the superpotential

$$\lambda_q^{ijk} \hat{\nu}_i^c \hat{q}_j \hat{q}_k^c . \quad (2.10)$$

Requiring that this term is allowed by the $U(1)_{extra}$, i.e. $Q_{\nu^c} = -Q_q - Q_{q^c}$, and using (2.7), we obtain that relation

$$Q_q + Q_{q^c} = Q_{H_1} + Q_{H_2} , \quad (2.11)$$

must be fulfilled. Taking into account this relation together with the equation of cancellation of the $[SU(3)_C]^2 - U(1)_{extra}$ anomaly, we finally obtain that the number of families of the exotic triplets must be $n_q = 3$.

Let us now consider the $[Gravity]^2 - U(1)_{extra}$ anomaly cancellation condition, that is, $\sum Q_{extra} = 0$, to obtain: $3(6Q_Q + 3Q_u + 3Q_d + 2Q_L + Q_e + Q_{\nu^c}) + 2n_H(Q_{H_1} + Q_{H_2}) + 3(3Q_q + 3Q_{q^c}) = 0$. Using (2.3-2.7) and (2.11), one finally gets $(2n_H - 3)(Q_{H_1} + Q_{H_2}) = 0$ which has no solution for $Q_{H_1} \neq -Q_{H_2}$ or an integer number of Higgs families. Then we conclude that we have to add more exotic matter to the spectrum in order to cancel the gravitational anomaly. Since we would like to extend the model with the minimal content of matter, the simplest solution is *to add extra singlets* under the Standard Model gauge group, in order not to alter the usual anomaly cancellation. In particular, we will add in principle n_s generations of singlets $\hat{s}(1, 1, 0, Q_s)$. Thus the anomaly cancellation condition implies that Q_s must have the value $Q_s = \frac{3-2n_H}{n_s}(Q_{H_1} + Q_{H_2})$.

Now, with the $[U(1)_Y]^2 - U(1)_{extra}$ anomaly cancellation condition, $\sum Y^2 Q_{extra} = 0$, where Y generically denotes the hypercharges of the particles, and using (2.3-2.6), (2.11) and the value for Q_d obtained above, one can find the following equation: $(9Y_q^2 + n_H - 4)(Q_{H_1} + Q_{H_2}) = 0$. Since we want to forbid the bilinear terms in the superpotential, we must impose $9Y_q^2 + n_H - 4 = 0$. For $n_H = 1, 2$ one obtains that Y_q must be an irrational number. The case $n_H = 3$ is excluded by the requirement of proton stability, as discussed above. For $n_H > 4$ we obtain a complex value for Y_q . Finally, with $n_H = 4$ the $[U(1)_{extra}]^2 - U(1)_Y$ anomaly cancellation condition, $\sum Q_{extra}^2 Y = 0$, implies $Q_{H_1} = -Q_{H_2}$, and the bilinear terms would be allowed in the superpotential.

We conclude with this analysis that it is not possible to cancel all the anomalies with only three new degrees of freedom (Q_q, Q_{q^c}, Q_s) . We have checked that neither is possible with four. In particular, we have considered the following possibilities: Two types of vector-like triplets, $\hat{q}_1(3, 1, Y_{q_1}, Q_{q_1})$, $\hat{q}_1^c(\bar{3}, 1, -Y_{q_1}, Q_{q_1^c})$, $\hat{q}_2(3, 1, Y_{q_2}, Q_{q_2})$ and $\hat{q}_2^c(\bar{3}, 1, -Y_{q_2}, Q_{q_2^c})$; one type of vector-like triplets, $\hat{q}(3, 1, Y_q, Q_q)$ and $\hat{q}^c(\bar{3}, 1, -Y_q, Q_{q^c})$, and two singlets with

opposite hypercharge for not altering the Standard Model anomaly cancellation, $\hat{s}_1(1, 1, Y_s, Q_{s_1})$ and $\hat{s}_2(1, 1, -Y_s, Q_{s_2})$; one type of vector-like triplets, $\hat{q}(3, 1, Y_q, Q_q)$ and $\hat{q}^c(\bar{3}, 1, -Y_q, Q_{q^c})$, one $SU(2)_L$ doublet with zero hypercharge, $\hat{l}(1, 2, 0, Q_l)$, and one Standard Model singlet $\hat{s}(1, 1, 0, Q_s)$; one type of vector-like triplets, $\hat{q}(3, 1, Y_q, Q_q)$ and $\hat{q}^c(\bar{3}, 1, -Y_q, Q_{q^c})$, and one type of vector-like doublets, $\hat{l}(1, 2, Y_l, Q_l)$ and $\hat{l}^c(1, 2, -Y_l, Q_{l^c})$. In all these cases it is not possible to cancel all the anomalies solving all the problems discussed above, and having effective mass terms for the exotic matter.

Let us then find the simplest model that cancels all anomalies. This model adds the following exotic matter to the spectrum (with five extra charges): three generations of vector-like triplets

$$\hat{q}(3, 1, Y_q, Q_q) , \quad \hat{q}^c(\bar{3}, 1, -Y_q, Q_{q^c}) , \quad (2.12)$$

and n_s generations of singlets

$$\hat{s}(1, 1, 0, Q_s) , \quad (2.13)$$

as obtained above, and in addition n_l generations of doublets

$$\hat{l}(1, 2, Y_l, Q_l) , \quad \hat{l}^c(1, 2, -Y_l, Q_{l^c}) . \quad (2.14)$$

The doublets must also be sufficiently massive as the extra triplets to have evaded the experimental detection. Then, we allow the following effective mass term in the superpotential:

$$\lambda_l^{ijk} \hat{\nu}_i^c \hat{l}_j^c \hat{l}_k^c . \quad (2.15)$$

As a consequence, the following condition is obtained:

$$Q_l + Q_{l^c} = Q_{H_1} + Q_{H_2} . \quad (2.16)$$

In addition, the $[SU(2)_L]^2 - U(1)_{extra}$ anomaly cancellation gives rise to:

$$Q_d = \frac{n_H + n_l - 6}{9} Q_{H_1} + \frac{n_H + n_l}{9} Q_{H_2} . \quad (2.17)$$

Note that the requirement (2.8) implies that $n_H + n_l \neq 3$.

From the $[Gravity]^2 - U(1)_{extra}$ anomaly cancellation we obtain the value of the extra charge of the singlet:

$$Q_s = \frac{3 - 2n_H - 2n_l}{n_s} (Q_{H_1} + Q_{H_2}) . \quad (2.18)$$

After replacing all the variables in the $[U(1)_Y]^2 - U(1)_{extra}$ anomaly cancellation condition we are left with the equation

$$18 Y_q^2 + 4 n_l Y_l^2 = 8 - n_l - 2 n_H . \quad (2.19)$$

The left side of the equation is a sum of positive quantities so we obtain an upper bound on the number of generations $n_l + 2n_H \leq 8$. We can study this equation searching for reasonable rational values of the hypercharges. The results are given in Table 1.

n_H	1	1	1	1	2	2	3	2	3
n_l	3	3	4	4	2	2	1	2	1
Y_q	$\pm\frac{2}{5}$	0	$\pm\frac{1}{3}$	$\pm\frac{1}{9}$	$\pm\frac{1}{5}$	$\pm\frac{1}{3}$	$\pm\frac{2}{9}$	0	0
Y_l	$\pm\frac{1}{10}$	$\pm\frac{1}{2}$	0	$\pm\frac{1}{3}$	$\pm\frac{2}{5}$	0	$\pm\frac{1}{6}$	$\pm\frac{1}{2}$	$\pm\frac{1}{2}$

Table 1: Number of generations of Higgses and extra doublets, and hypercharges that solve the $[U(1)_Y]^2 - U(1)_{extra}$ anomaly equation.

The equation associated to $[U(1)_{extra}]^2 - U(1)_Y$ is quadratic in the extra charges:

$$3(6Y_Q Q_Q^2 + 3Y_u Q_u^2 + 3Y_d Q_d^2 + 2Y_L Q_L^2 + Y_e Q_e^2 + Y_\nu Q_{\nu^c}^2) + n_H(2Y_{H_1} Q_{H_1}^2 + 2Y_{H_2} Q_{H_2}^2) + n_q(3Y_q Q_q^2 - 3Y_q Q_{q^c}^2) + n_l(2Y_l Q_l^2 - 2Y_l Q_{l^c}^2) + n_s Y_s Q_s^2 = 0 . \quad (2.20)$$

After substituting all the extra charges known, we obtain the value of Q_q in terms of Q_{H_1} , Q_{H_2} , and Q_l . Using (2.11) we also obtain Q_{q^c} in terms of Q_{H_1} , Q_{H_2} , and Q_l .

Finally, the equation associated to $[U(1)_{extra}]^3$ is cubic in the extra charges, $\sum Q_{extra}^3 = 0$. We study this equation for each set of (n_H, n_l, Y_q, Y_l) of Table 1. We obtain Q_l in terms of Q_{H_1} and Q_{H_2} . Using (2.16), we also obtain Q_{l^c} in terms of Q_{H_1} and Q_{H_2} . The only cases that give rise to rational values for Q_l and Q_{l^c} have the following number of generations:

$$n_H = 1 , \quad n_l = 3 , \quad n_s = 6 , \quad (2.21)$$

with

$$Y_q = \pm\frac{2}{5} , \quad Y_l = \frac{1}{10} , \quad (2.22)$$

and two distinct solutions for Q_l as shown in the right side of Table 2, or

$$Y_q = 0 , \quad Y_l = \frac{1}{2} , \quad (2.23)$$

and two distinct solutions for Q_q , as shown also in the right side of Table 2.

It is worth noticing here that, although at the end we are left with the six different solutions (models) discussed above, we will see in the next section that all of them give rise to the same phenomenology at low energies. This is because the six models only differ in the extra charges and hypercharges of the exotic matter, and this matter does not play any role in the EW breaking.

We have then obtained all the extra charges in terms of two of them, Q_{H_1} and Q_{H_2} . For rational values of Q_{H_1} and Q_{H_2} we obtain rational values for the rest of extra charges. For definiteness, we add two additional conditions for the complete determination of the extra charges. First, we impose that the bases of the hypercharge Y and the extra charge Q are orthogonal, i.e. $Tr[YQ] = 0$. This implies $Q_{H_2} = 6 Q_{H_1}$. Second, following [26] we impose the normalization condition for the extra charges $Tr[Y^2] = Tr[Q^2]$. This condition is in fact non physical since the relevant quantity is the product of the extra gauge coupling constant g'_1 by the normalization factor. From this condition we obtain the value of Q_{H_1} and, consequently, the values of all the extra charges. We show these values in Tables 2 and 3, where the normalization factor is given by $N = \sqrt{\frac{3}{2426}}$.

Q_q	Q_{q^c}	Q_l	Q_{l^c}	
$\frac{257}{30}N$	$\frac{373}{30}N$	$\frac{19}{15}N$	$\frac{296}{15}N$	Model 1: $Y_q = \frac{2}{5}$, $Y_l = \frac{1}{10}$, $Q_l = \frac{1}{45}(-5Q_{H_1} + 4Q_{H_2})$
$\frac{173}{30}N$	$\frac{457}{30}N$	$\frac{271}{15}N$	$\frac{44}{15}N$	Model 2: $Y_q = \frac{2}{5}$, $Y_l = \frac{1}{10}$, $Q_l = \frac{1}{45}(31Q_{H_1} + 40Q_{H_2})$
$\frac{373}{30}N$	$\frac{257}{30}N$	$\frac{19}{15}N$	$\frac{296}{15}N$	Model 3: $Y_q = -\frac{2}{5}$, $Y_l = \frac{1}{10}$, $Q_l = \frac{1}{45}(-5Q_{H_1} + 4Q_{H_2})$
$\frac{457}{30}N$	$\frac{173}{30}N$	$\frac{271}{15}N$	$\frac{44}{15}N$	Model 4: $Y_q = -\frac{2}{5}$, $Y_l = \frac{1}{10}$, $Q_l = \frac{1}{45}(31Q_{H_1} + 40Q_{H_2})$
$\frac{7}{2}N$	$\frac{35}{2}N$	$\frac{19}{3}N$	$\frac{44}{3}N$	Model 5: $Y_q = 0$, $Y_l = \frac{1}{2}$, $Q_q = \frac{1}{6}(Q_{H_1} + Q_{H_2})$
$\frac{35}{2}N$	$\frac{7}{2}N$	$\frac{19}{3}N$	$\frac{44}{3}N$	Model 6: $Y_q = 0$, $Y_l = \frac{1}{2}$, $Q_q = \frac{5}{6}(Q_{H_1} + Q_{H_2})$

Table 2: Values of the $U(1)_{extra}$ charges of the extra triplets and doublets added to the Standard Model spectrum of the $\mu\nu$ SSM, for the six solutions of the $[U(1)_{extra}]^3$ anomaly cancellation condition.

$Q_{H_1} = 3N$	$Q_{H_2} = 18N$	$Q_Q = -\frac{31}{3}N$	$Q_u = -\frac{23}{3}N$	$Q_d = \frac{22}{3}N$
$Q_L = 3N$	$Q_e = -6N$	$Q_{\nu^c} = -21N$	$Q_s = -\frac{35}{2}N$	

Table 3: Values of the $U(1)_{extra}$ charges for the Standard Model content of the $\mu\nu$ SSM and for the extra singlets.

Summarizing, we have found six interesting models with the following exotic matter: three generations of vector-like color triplets with respect to the Standard Model gauge group (2.12), three generations of $SU(2)_L$ doublets (2.14), and six¹ Standard Model singlets (2.13). The superpotential is given by:

$$\begin{aligned}
W = & \epsilon_{ab}(Y_u^{ij}\hat{H}_2^a\hat{Q}_i^b\hat{u}_j^c + Y_d^{ij}\hat{H}_1^a\hat{Q}_i^b\hat{d}_j^c + Y_e^{ij}\hat{H}_1^a\hat{L}_i^b\hat{e}_j^c + Y_\nu^{ij}\hat{H}_2^b\hat{L}_i^a\hat{\nu}_j^c) - \epsilon_{ab}\lambda^i\hat{\nu}_i^c\hat{H}_1^a\hat{H}_2^b \\
& + \epsilon_{ab}(\lambda^{ijk}\hat{Q}_i^a\hat{L}_j^b\hat{d}_k^c + \lambda^{\prime\prime\prime ijk}\hat{L}_i^a\hat{L}_j^b\hat{e}_k^c) + \lambda_q^{ijk}\hat{\nu}_i^c\hat{q}_j\hat{q}_k^c + \epsilon_{ab}\lambda_l^{ijk}\hat{\nu}_i^c\hat{l}_j^a(\hat{l}_k^b)^c. \quad (2.24)
\end{aligned}$$

The extra singlets do not have couplings in the superpotential. The extra triplets and doublets have effective mass terms in the superpotential avoiding conflicts with the experimental searches of exotic matter. The singlets, as they do not couple in the superpotential, and only interact through their extra $U(1)$ charge, do not need to be massive for escaping detection since the value of the lower bound on the mass of an extra Z is quite large.

Let us now make a comment on the hypercharges of the extra matter. In these models, the hypercharges of the exotic matter lead to non-standard fractional electric charges. This issue has been discussed for example in [27], and references therein. In the case of extra triplets, they could form color-neutral fractionally charged states since the triplets can bind. The lightest of these states will be stable due to electric charge conservation. As pointed out in [28], the estimation of its relic abundance contradicts limits on the existence of fractional charge in matter which is less than 10^{-20} per nucleon [29]. Thus, avoiding such fractionally charged states is necessary. A possible mechanism to carry it out is

¹If preferred, one could imagine that those six generations are in fact $3 + 3$ generations that could be distinguished by some high-energy extra $U(1)$ gauge group, perhaps coming from the compactification of a string model [20].

inflation. Inflation would dilute these particles. The reheating temperature T_{RH} should be low enough not to produce them again. This reheating temperature must be smaller than 10^{-3} times the mass of the particle [30], so in our case $T_{RH} < 1$ GeV. This, in principle is possible since the only constraint on this temperature is to be larger than 1 MeV not to spoil the successful nucleosynthesis predictions.

Finally, let us recall that the models where R -parity is conserved still need some fine-tuning to agree with the experimental bounds on the proton lifetime. This is because R -parity does not forbid non-renormalizable dimension five operators that break baryon or lepton number, and could produce too fast proton decay if the couplings are of order one [31]. We have checked that although in the model analyzed here, there are 43 non-renormalizable dimension five baryon number violating operators allowed by the gauge symmetry of the Standard Model, such as for example $\hat{Q}\hat{Q}\hat{Q}\hat{L}$, $\hat{u}^c\hat{u}^c\hat{d}^c\hat{e}^c$ or $\hat{Q}\hat{Q}\hat{Q}\hat{H}_1$, all of them turn out to be forbidden by the extra $U(1)$.

3. Electroweak breaking and experimental constraints

The gauge symmetry discussed in the previous section $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{extra}$ has to be spontaneously broken to $SU(3)_C \times U(1)_{e.m.}$. To discuss this breaking we have first to calculate the neutral scalar potential, which is the sum of three contributions: F-terms, D-terms and soft terms. Working in the framework of gravity-mediated supersymmetry breaking, and taking into account the superpotential (2.24), the latter are given by:

$$\begin{aligned} \mathcal{L}_{soft} = & \frac{1}{2}(M_3\tilde{\lambda}_3\tilde{\lambda}_3 + M_2\tilde{\lambda}_2\tilde{\lambda}_2 + M_1\tilde{\lambda}_1\tilde{\lambda}_1 + M'_1\tilde{\lambda}'_1\tilde{\lambda}'_1 + h.c.) \\ & - \epsilon_{ab}[(A_u Y_u)^{ij} H_2^b \tilde{Q}_i^a \tilde{u}_j^c + (A_d Y_d)^{ij} H_1^b \tilde{Q}_i^a \tilde{d}_j^c + (A_e Y_e)^{ij} H_1^b \tilde{L}_i^a \tilde{e}_j^c + (A_\nu Y_\nu)^{ij} H_2^b \tilde{L}_i^a \tilde{\nu}_j^c \\ & + (A_{\lambda'} \lambda')^{ijk} \tilde{Q}_i^a \tilde{L}_j^b \tilde{d}_k^c + (A_{\lambda''} \lambda'')^{ijk} \tilde{L}_i^a \tilde{L}_j^b \tilde{e}_k^c - (A_\lambda \lambda)^i \tilde{\nu}_i^c H_1^a H_2^b + (A_{\lambda_l} \lambda_l)^{ijk} \tilde{\nu}_i^c \tilde{l}_j^a \tilde{l}_b^{kc} + h.c.] \\ & - [(A_{\lambda_q} \lambda_q)^{ijk} \tilde{\nu}_i^c \tilde{q}_j^a \tilde{q}_k^c + h.c.] - [(M_{\tilde{Q}}^2)^{ij} \tilde{Q}_i^a \tilde{Q}_j^a + (M_{\tilde{u}^c}^2)^{ij} \tilde{u}_i^c \tilde{u}_j^c + (M_{\tilde{d}^c}^2)^{ij} \tilde{d}_i^c \tilde{d}_j^c \\ & + (M_{\tilde{L}}^2)^{ij} \tilde{L}_i^a \tilde{L}_j^a + (M_{\tilde{e}^c}^2)^{ij} \tilde{e}_i^c \tilde{e}_j^c + M_{H_1}^2 H_1^{a*} H_1^a + M_{H_2}^2 H_2^{a*} H_2^a + (M_{\tilde{\nu}^c}^2)^{ij} \tilde{\nu}_i^c \tilde{\nu}_j^c \\ & + (M_{\tilde{s}}^2)^{ij} \tilde{s}_i^* \tilde{s}_j + (M_{\tilde{q}}^2)^{ij} \tilde{q}_i^* \tilde{q}_j + (M_{\tilde{q}^c}^2)^{ij} \tilde{q}_i^{c*} \tilde{q}_j^c + (M_{\tilde{l}}^2)^{ij} \tilde{l}_i^* \tilde{l}_j + (M_{\tilde{l}^c}^2)^{ij} \tilde{l}_i^{ac*} \tilde{l}_j^{ac}] . \end{aligned} \quad (3.1)$$

Once the EW symmetry is spontaneously broken, the neutral scalars develop in general the following VEVs:

$$\langle H_1^0 \rangle = v_1, \quad \langle H_2^0 \rangle = v_2, \quad \langle \tilde{\nu}_i \rangle = \nu_i \quad \langle \tilde{\nu}_i^c \rangle = \nu_i^c . \quad (3.2)$$

We have checked that the neutral components of the exotic matter do not take VEVs in a wide region of the parameter space, where we will concentrate. In what follows, it will be enough for our purposes to neglect mixing between generations in (2.24) and (3.1), and to assume that only one generation of sneutrinos gets VEVs, ν , ν^c . The extension of the analysis to all generations is straightforward, and the conclusions are similar. The

expression of the neutral scalar potential is then given by:

$$\begin{aligned}
\langle V^0 \rangle = & \frac{1}{8}(g_1^2 + g_2^2)(|v_1|^2 + |\nu|^2 - |v_2|^2)^2 \\
& + \frac{1}{2}g_1'^2(Q_{H_1}|v_1|^2 + Q_{H_2}|v_2|^2 + Q_L|\nu|^2 + Q_{\nu^c}|\nu^c|^2)^2 \\
& + |Y_\nu|^2(|v_2|^2|\nu^c|^2 + |v_2|^2|\nu|^2 + |\nu|^2|\nu^c|^2) \\
& + |\lambda|^2(|v_1|^2|v_2|^2 + |\nu^c|^2|v_2|^2 + |\nu^c|^2|v_1|^2) \\
& + (-\lambda Y_\nu^* v_1 \nu^* |v_2|^2 - \lambda Y_\nu^* v_1 \nu^* |\nu^c|^2 + h.c.) \\
& + M_L^2|\nu|^2 + M_{\nu^c}^2|\nu^c|^2 + M_{H_1}^2|v_1|^2 + M_{H_2}^2|v_2|^2 \\
& + (A_\nu Y_\nu v_2 \nu \nu^c - A_\lambda \lambda \nu^c v_1 v_2 + h.c.) .
\end{aligned} \tag{3.3}$$

We also assume, for simplicity, that there is not CP violation in the scalar sector and we take all the parameters and VEVs real in (3.3). The four minimization conditions with respect to the VEVs v_1 , v_2 , ν^c , ν , are:

$$\begin{aligned}
& \frac{1}{4}(g_1^2 + g_2^2)(v_1^2 + \nu^2 - v_2^2)v_1 + g_1'^2(Q_{H_1}v_1^2 + Q_{H_2}v_2^2 + Q_L\nu^2 + Q_{\nu^c}\nu^{c2})Q_{H_1}v_1 \\
& + \lambda^2 v_1(v_2^2 + \nu^{c2}) + M_{H_1}^2 v_1 - \lambda Y_\nu \nu |v_2|^2 - \lambda Y_\nu \nu |\nu^c|^2 - A_\lambda \lambda \nu^c v_2 = 0 , \\
& - \frac{1}{4}(g_1^2 + g_2^2)(v_1^2 + \nu^2 - v_2^2)v_2 + g_1'^2(Q_{H_1}v_1^2 + Q_{H_2}v_2^2 + Q_L\nu^2 + Q_{\nu^c}\nu^{c2})Q_{H_2}v_2 \\
& + Y_\nu^2 v_2(\nu^2 + \nu^{c2}) + \lambda^2 v_2(v_1^2 + \nu^{c2}) + M_{H_2}^2 v_2 - 2\lambda Y_\nu v_1 \nu v_2 \\
& + A_\nu Y_\nu \nu \nu^c - A_\lambda \lambda \nu^c v_1 = 0 , \\
& g_1'^2(Q_{H_1}v_1^2 + Q_{H_2}v_2^2 + Q_L\nu^2 + Q_{\nu^c}\nu^{c2})Q_{\nu^c}\nu^c + Y_\nu^2 \nu^c(v_2^2 + \nu^2) - A_\lambda \lambda v_1 v_2 \\
& + \lambda^2 \nu^c(v_1^2 + v_2^2) + M_{\nu^c}^2 \nu^c - 2\lambda Y_\nu v_1 \nu \nu^c + A_\nu Y_\nu v_2 \nu = 0 , \\
& \frac{1}{4}(g_1^2 + g_2^2)(v_1^2 + \nu^2 - v_2^2)\nu + g_1'^2(Q_{H_1}v_1^2 + Q_{H_2}v_2^2 + Q_L\nu^2 + Q_{\nu^c}\nu^{c2})Q_L\nu \\
& + Y_\nu^2 \nu(v_2^2 + \nu^{c2}) + M_L^2 \nu - \lambda Y_\nu v_1 v_2^2 - \lambda Y_\nu v_1 \nu^{c2} + A_\nu Y_\nu v_2 \nu^c = 0 .
\end{aligned} \tag{3.4}$$

Notice that in the last equation in (3.4) $\nu \rightarrow 0$ as $Y_\nu \rightarrow 0$, and since the coupling Y_ν determines the Dirac mass for the neutrinos, $m_D \equiv Y_\nu v_2$, then $Y_\nu \sim 10^{-6}$, and therefore ν has to be very small. The smallness of the left-handed sneutrino VEVs for a correct description of the neutrino sector in the $\mu\nu$ SSM, compatible with current data, has been proved in [9, 11, 8].

We can now approximate the minimization equations neglecting the values of ν and Y_ν , and we are left with only three equations. Solving the minimization conditions for the soft masses in terms of the extra charges, coupling constants, VEVs, and the parameters

λ and $A_\lambda \lambda$, one obtains:

$$\begin{aligned}
M_{H_1}^2 &= -\frac{1}{4}(g_1^2 + g_2^2)(v_1^2 - v_2^2) - g_1'^2(Q_{H_1}v_1^2 + Q_{H_2}v_2^2 + Q_{\nu^c}\nu^{c2})Q_{H_1} \\
&\quad - \lambda^2(v_2^2 + \nu^{c2}) + A_\lambda \lambda \nu^c \frac{v_2}{v_1} , \\
M_{H_2}^2 &= \frac{1}{4}(g_1^2 + g_2^2)(v_1^2 - v_2^2) - g_1'^2(Q_{H_1}v_1^2 + Q_{H_2}v_2^2 + Q_{\nu^c}\nu^{c2})Q_{H_2} \\
&\quad - \lambda^2(v_1^2 + \nu^{c2}) + A_\lambda \lambda \nu^c \frac{v_1}{v_2} , \\
M_{\nu^c}^2 &= -g_1'^2(Q_{H_1}v_1^2 + Q_{H_2}v_2^2 + Q_{\nu^c}\nu^{c2})Q_{\nu^c} - \lambda^2(v_1^2 + v_2^2) + A_\lambda \lambda \frac{v_1 v_2}{\nu^c} . \quad (3.5)
\end{aligned}$$

Note that these equations are equivalent (substituting ν^c by the VEV of a singlet scalar) to the minimization conditions for the $U(1)$ SSM models [32, 26], where correct EW breaking is known to take place.

On the other hand, the VEVs have to satisfy several phenomenological constraints. First, the mass of the W boson, $M_W = \frac{1}{2}g_2^2(v_1^2 + v_2^2 + \nu^2)$, is well determined, leading to $(v_1^2 + v_2^2) \simeq (174 \text{ GeV})^2$ when ν is neglected. Second, the Z boson of the Standard Model and the Z' boson associated to the $U(1)_{extra}$ are mixed with a mass-squared matrix given by:

$$\begin{pmatrix} M_Z^2 & M_{ZZ'}^2 \\ M_{ZZ'}^2 & M_{Z'}^2 \end{pmatrix} , \quad (3.6)$$

where the entries are functions of the VEVs, gauge coupling constants and extra charges,

$$\begin{aligned}
M_Z^2 &= \frac{1}{2}(g_1^2 + g_2^2)(v_1^2 + v_2^2) , \\
M_{Z'}^2 &= 2g_1'^2(Q_{H_1}^2 v_1^2 + Q_{H_2}^2 v_2^2 + Q_{\nu^c}^2 \nu^{c2}) , \\
M_{ZZ'}^2 &= g_1' \sqrt{g_1^2 + g_2^2} (-Q_{H_1} v_1^2 + Q_{H_2} v_2^2) . \quad (3.7)
\end{aligned}$$

Diagonalizing this matrix one obtains the mass eigenstates. The experimental constraints imply the following bound [33] for the mixing parameter

$$R = \frac{(M_{ZZ'}^2)^2}{M_Z^2 M_{Z'}^2} \leq 10^{-3} . \quad (3.8)$$

In addition, the mass of the heaviest eigenstate should be larger than about 1 TeV [34]. If we also ask the heaviest eigenstate to be lighter than 2000 GeV in order not to have a very large fine-tuning (and for the Z' to be discovered at present accelerator experiments), then

$$\begin{aligned}
(1000)^2 &\leq \frac{1}{4}(g_1^2 + g_2^2)(v_1^2 + v_2^2) + g_1'^2(Q_{H_1}^2 v_1^2 + Q_{H_2}^2 v_2^2 + Q_{\nu^c}^2 \nu^{c2}) \\
&\quad + [\frac{1}{16}(g_1^2 + g_2^2)^2(v_1^2 + v_2^2)^2 + g_1'^4(Q_{H_1}^2 v_1^2 + Q_{H_2}^2 v_2^2 + Q_{\nu^c}^2 \nu^{c2})^2 \\
&\quad - \frac{1}{2}(g_1^2 + g_2^2)g_1'^2(v_1^2 + v_2^2)(Q_{H_1}^2 v_1^2 + Q_{H_2}^2 v_2^2 + Q_{\nu^c}^2 \nu^{c2}) \\
&\quad + g_1'^2(g_1^2 + g_2^2)(-Q_{H_1} v_1^2 + Q_{H_2} v_2^2)^2]^{1/2} \leq (2000)^2 . \quad (3.9)
\end{aligned}$$

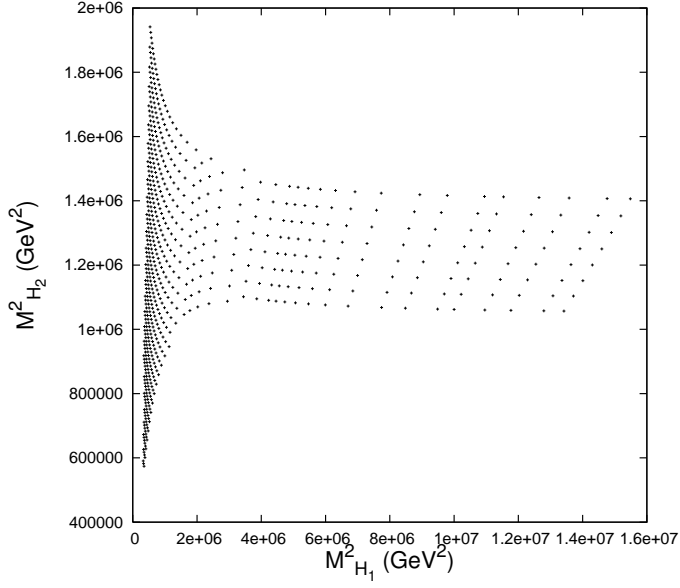


Figure 1: Allowed region by the experimental constraints on the Z' in the plane $M_{H_1}^2 - M_{H_2}^2$, for $\lambda = 0.1$

From the above equations it is obvious that the six models found in the previous section give rise to the same phenomenology at low energies, since they only differ in the extra charges and hypercharges of the exotic matter, and this matter does not play any role in the EW breaking.

In order to study the solutions of the equations, we assume the following reasonable values for the parameters: $A_\lambda \lambda = 0.1$ TeV and $\lambda = 0.1, 0.3$. For the sake of definiteness we also take $g'_1 = g_1$, $1 < \tan \beta = \frac{v_2}{v_1} < 35$, and we work in the parameter space $(\nu^c, \tan \beta)$. Once imposed the experimental constraints on the existence of a new gauge boson Z' , we have checked that the effect of the bound on the $Z - Z'$ mixing is more important than the bounds on the mass of the heaviest eigenstate, although it is still possible to find wide allowed regions. The former experimental constraint implies a lower bound on the VEV of the right-handed sneutrino ν^c , depending on the value of $\tan \beta$. In particular, for $\lambda = 0.3$ and $\tan \beta = 1$, ν^c must be larger than 2 TeV. For increasing values of $\tan \beta$, the lower bound on ν^c increases since it is more difficult to suppress the $Z - Z'$ mixing. For example, for $\tan \beta = 3$ (7), one obtains that ν^c must be larger than about 4 (4.6) TeV. For $\tan \beta$ larger than 7, the lower bound on ν^c practically does not vary. Similar results are obtained for $\lambda = 0.1$, although in this case a tachyonic region appears and we always need values of ν^c larger than 2.5 TeV.

One can translate the constraints on the Z' to the plane $(M_{H_1}^2, M_{H_2}^2)$, finding the allowed region in the parameter space of the soft masses. We show these regions in Figs. 1 and 2 for $\lambda = 0.1$ and 0.3, respectively.

Once we have shown that the model is in principle phenomenologically viable, let us now focus our attention on the neutralino sector. In the $\mu\nu$ SSM with an extra $U(1)$ gauge symmetry, the MSSM neutralinos mix with the extra gaugino. The fact that R -parity is

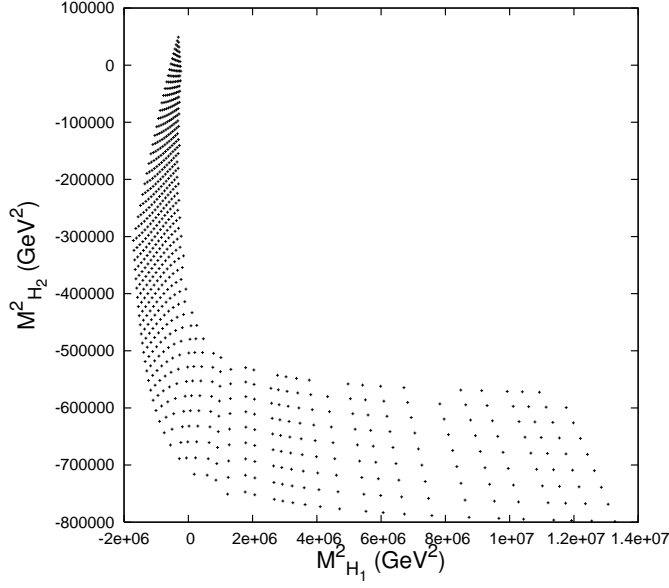


Figure 2: Allowed region by the experimental constraints on the Z' in the plane $M_{H_1}^2 - M_{H_2}^2$, for $\lambda = 0.3$

broken in this model, also produces the mixing of the neutralinos with the left- and right-handed neutrinos. Of course, now we have to be sure that one eigenvalue of this matrix is very small, reproducing the experimental results about neutrino masses. In the weak interaction basis defined by $\psi^{0t} = (\tilde{Z}', \tilde{B}^0 = -i\tilde{\lambda}', \tilde{W}_3^0 = -i\tilde{\lambda}_3, \tilde{H}_1^0, \tilde{H}_2^0, \nu^c, \nu)$, the neutral fermion mass terms in the Lagrangian are $\mathcal{L}_{neutral}^{mass} = -\frac{1}{2}(\psi^0)^t \mathcal{M}_n \psi^0 + h.c.$, with \mathcal{M}_n a 7×7 (11×11 if we include all generations of neutrinos) matrix,

$$\mathcal{M}_n = \begin{pmatrix} M & m \\ m^t & 0 \end{pmatrix}, \quad (3.10)$$

where

$$M = \begin{pmatrix} M'_1 & 0 & 0 & \sqrt{2}g'_1 Q_{H_1} v_1 & \sqrt{2}g'_1 Q_{H_2} v_2 & \sqrt{2}g'_1 Q_{\nu^c} \nu^c \\ 0 & M_1 & 0 & -\frac{1}{\sqrt{2}}g_1 v_1 & \frac{1}{\sqrt{2}}g_1 v_2 & 0 \\ 0 & 0 & M_2 & \frac{1}{\sqrt{2}}g_2 v_1 & -\frac{1}{\sqrt{2}}g_2 v_2 & 0 \\ \sqrt{2}g'_1 Q_{H_1} v_1 & -\frac{1}{\sqrt{2}}g_1 v_1 & \frac{1}{\sqrt{2}}g_2 v_1 & 0 & -\lambda \nu^c & -\lambda v_2 \\ \sqrt{2}g'_1 Q_{H_2} v_2 & \frac{1}{\sqrt{2}}g_1 v_2 & -\frac{1}{\sqrt{2}}g_2 v_2 & -\lambda \nu^c & 0 & -\lambda v_1 + Y_\nu \nu \\ \sqrt{2}g'_1 Q_{\nu^c} \nu^c & 0 & 0 & -\lambda v_2 & -\lambda v_1 + Y_\nu \nu & 0 \end{pmatrix}, \quad (3.11)$$

is very similar to the neutralino mass matrix of the $U(1)$ SSM (substituting ν^c by the VEV of a singlet scalar and neglecting the contributions $Y_\nu \nu$), and

$$m^t = \left(\sqrt{2}g'_1 Q_\nu \nu \quad -\frac{1}{\sqrt{2}}g_1 \nu \quad \frac{1}{\sqrt{2}}g_2 \nu \quad 0 \quad Y_\nu \nu^c \quad Y_\nu v_2 \right). \quad (3.12)$$

Using typical values of the soft gaugino masses, and with values for the rest of parameters in the region allowed by the constraints on the Z' , we have checked numerically that correct neutrino masses can easily be obtained, i.e. once we diagonalize the neutralino mass matrix, one eigenvalue is sufficiently small, of the order of 10^{-2} eV.

However, for the general case of three generations of left- and right-handed neutrinos, unlike what occurs for the $\mu\nu$ SSM [8], the analysis is not so straightforward. As discussed in Section 2, the presence of the extra $U(1)$ group forbids the usual Majorana mass term, $\kappa\hat{\nu}^c\hat{\nu}^c\hat{\nu}^c$, of the original $\mu\nu$ SSM (1.1). Thus, taking into account the generalization of (3.11) for three generations [7], right-handed neutrinos can only acquire large masses through the mixings with the extra gaugino and the Higgsinos due to the terms proportional to $g'_1\nu_i^c$ and λ^i , respectively. Without loss of generality, one can define a basis ν_i^c such that the VEVs of the right-handed sneutrinos are in the ν_1^c direction ($\nu_{2,3}^c = 0$), while $\lambda^3 = 0$. Then, only two right-handed neutrinos, ν_1^c and ν_2^c can have EW-scale masses, denoted as $M_{\nu_1^c}$ and $M_{\nu_2^c}$. The third one combines with the left-handed neutrinos to form a nearly massless Dirac particle. As a consequence, the EW-scale see-saw only works for two linear combinations of left-handed neutrinos, that we will denote as ν_1 and ν_2 , with a mass given by² $m_{\nu_i} \sim \frac{(Y_\nu^{ii}v_2)^2}{M_{\nu_i^c}}$ for $i = 1, 2$, with $Y_\nu \sim 10^{-6}$. In general, the four light neutrino states ($\nu_1, \nu_2, \nu_3, \nu_3^c$) are mixed with the following mass matrix,

$$\begin{pmatrix} m_{\nu_1} & 0 & 0 & m_D^1 \\ 0 & m_{\nu_2} & 0 & m_D^2 \\ 0 & 0 & 0 & m_D^3 \\ m_D^1 & m_D^2 & m_D^3 & 0 \end{pmatrix}, \quad (3.13)$$

where the state ν_3 is orthogonal to $\nu_{1,2}$ states and $m_D^k \simeq Y_\nu^{k3}v_2$, with $k = 1, 2, 3$, are the Dirac masses in this basis. At tree-level, there are four light Majorana states. To account for neutrino data, for example the cosmological bound from WMAP on the sum of all light neutrino masses, of the order of the eV, a fine-tuning on the other three entries of the Yukawa matrix is necessary, $Y_\nu^{k3} \leq 10^{-11}$. We have checked this result numerically. For a more detailed discussion of this mechanism see [35], where a similar situation occurs in the context of $U(1)_{B-L}$ supersymmetric models with broken R-parity, although in that case the EW-scale see-saw works only for one neutrino.

Therefore, the $\mu\nu$ SSM with an extra $U(1)$, in the general case of three generations of neutrinos, predicts the existence of two heavy right-handed neutrinos, of the order of the TeV, and four light (three active and one sterile) neutrinos. Given the oscillation anomalies in LSND, MiniBooNE and MINOS [36, 37, 38], the extra light sterile neutrino might be welcome [39].

Finally, we have also performed an estimation of the tree-level upper bound on the lightest Higgs mass in this model. Let us recall that, neglecting the small neutrino Yukawa coupling effects, the expression of the upper bound on the lightest Higgs mass in the $\mu\nu$ SSM is equivalent to that of the NMSSM, once we define $\lambda^2 = \lambda_1^2 + \lambda_2^2 + \lambda_3^2$, and this is given

²Note that there are other contributions to neutrino masses due to the mixing of left-handed neutrinos with Higgsinos and neutral gauginos [8].

by the following tree-level expression [7]:

$$m_h^2 \leq M_Z^2 (\cos^2 2\beta + \frac{2\lambda^2 \cos^2 \theta_W}{g_2^2} \sin^2 2\beta) . \quad (3.14)$$

This bound receives a positive contribution from the extra $U(1)$ sector [40], in such a way that the tree-level formula for the upper bound on the lightest Higgs mass in the $\mu\nu$ SSM with an extra $U(1)$ is given by:

$$m_h^2 \leq M_Z^2 (\cos^2 2\beta + \frac{2\lambda^2 \cos^2 \theta_W}{g_2^2} \sin^2 2\beta) + 2g_1' v^2 (Q_{H_2} \cos^2 \beta + Q_{H_1} \sin^2 \beta)^2 . \quad (3.15)$$

It is worth recalling here that the LEP lower bound on the Higgs mass is 114 GeV. In the case of the MSSM, the tree-level upper bound on the lightest Higgs mass is M_Z , i.e. smaller than 114 GeV. As is well known, the MSSM is not still ruled out by LEP because the radiative corrections can raise this upper bound, although this can require some fine-tuning. In the case of the NMSSM, once perturbativity is imposed, the tree-level upper bound on the lightest Higgs mass can be as high as 110 GeV, improving the Higgs-mass problem of the MSSM. In [7], this issue has been analyzed in the context of the $\mu\nu$ SSM, and the upper bound turns out to be similar to the one of the NMSSM.

In the $\mu\nu$ SSM with an extra $U(1)$, the upper bound depends on the value of the extra gauge coupling g_1' , as shown in (3.15). Whereas for $g_1' \simeq g_1$ this bound is only raised to 113 GeV, for $g_1' \simeq 2g_1$ it is raised to about 120 GeV. Thus, the addition of an extra $U(1)$ gauge group to the $\mu\nu$ SSM has also the nice feature of increasing the tree-level bound on the Higgs mass leading to a larger window for the discovery of the Higgs at collider experiments.

4. Conclusions

The $\mu\nu$ SSM solves the μ problem of the MSSM and generates correct neutrino masses by simply using right-handed neutrinos. This mechanism implies that only dimensionless trilinear terms, breaking R -parity, are present in the superpotential. The non-presence in the superpotential of proton decay operators breaking R -parity, a trilinear term generating a domain wall problem, and bilinear terms such as the μ term and the Majorana masses, is solved in the $\mu\nu$ SSM using string theory arguments, discrete symmetries or non-renormalizable operators. In this work we have used a different strategy, namely an extra $U(1)$ gauge symmetry is added to the gauge group of the Standard Model. Since all fields of the $\mu\nu$ SSM can be charged under the extra $U(1)$, all the dangerous operators mentioned above could in principle be forbidden. We have checked that this is precisely the case. For example, dimension four and five baryon number violating operators are forbidden in the superpotential, ensuring the stability of the proton.

On the other hand, the anomaly cancellation conditions associated to the extra $U(1)$, allow us to constrain the values of the $U(1)$ charges. We have checked that six assignments of the $U(1)$ charges to the matter fields are viable, once extra matter is introduced. In

particular, three generations of vector-like color triplets and $SU(2)_L$ doublets, as well as six Standard Model singlets are needed.

We have studied the phenomenology of the model focusing our attention on the electroweak symmetry breaking. We have found that it is viable, with wide regions of the parameter space fulfilling the experimental constraints on the existence of a new gauge boson Z' . We have also studied the neutralino sector of the model, since the neutrinos and the extra gaugino mix with the MSSM neutralinos. We have checked numerically that the experimental results on neutrino masses can be reproduced. Nevertheless, the analysis of the generalized see-saw matrix in the case of three generations is different from the usual one in the $\mu\nu$ SSM. This is because of the absence of the trilinear term generating the domain wall problem, but also an effective Majorana mass term. Due to the absence of the latter, the EW-scale see-saw works only for two neutrinos and therefore one needs some entries of the Yukawa matrix small. Finally, we have estimated the tree-level upper bound on the lightest Higgs mass, finding that it can be as large as about 120 GeV.

Acknowledgments

This work was supported in part by the MICINN under grants FPA2009-08958 and FPA2009-09017, by the Comunidad de Madrid under grant HEPHACOS S2009/ESP-1473, and by the European Union under the Marie Curie-ITN program PITN-GA-2009-237920. The authors also acknowledge the support of the Consolider-Ingenio 2010 Programme under grant MultiDark CSD2009-00064.

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